Systematics of α-nucleus optical potentials

P. Mohr, H. Abele, U. Atzrott, G. Staudt (Physikalisches Institut, Univ. Tübingen),
R. Bieber, K. Grün, H. Oberhummer (Institut für Kernphysik, TU Wien),
T. Rauscher (Institut für Kernchemie, Univ. Mainz)
E. Somorjai (ATOMKI, Debrecen)

1. Introduction

For the description of nuclear processes in many astrophysical scenarios the knowledge of α-nucleus potentials is necessary. Such processes are radiative capture, transfer reactions and alpha-decay occuring in primordial nucleosynthesis, stellar hydrostatic and explosive burning modes.

Until now such nuclear processes have often been described using mainly phenomenological and energy-independent potentials (e.g. square well, Woods-Saxon potentials etc.). In this work we develop α-nucleus potentials using the folding procedure [1,2]. With this method the ambiguity of the phenomenological potentials can be avoided to a great extent. The uniqueness and the energy dependence of these potentials are an important feature with respect to astrophysical applications.

Such α-nucleus potentials have been used successfully for the description of scattering processes [3,4,5], transfer reactions [6,7,8] and radiative capture [9,10] on light nuclei. In this work we extend our systematic investigation of α-nucleus potentials to intermediate and heavy stable and unstable nuclei ($A \geq 70$) which are relevant for the p-process [11].

2. Folding procedure

The real part of the optical potential is deduced in the framework of the double-folding model of Kobos et al. [1] and is described by

$$V(r) = \lambda \int d\vec{r}_1 \int d\vec{r}_2 \rho_T(\vec{r}_1)\rho_\alpha(\vec{r}_2)t(E,\rho_T,\rho_\alpha,\vec{s}=\vec{r}+\vec{r}_2-\vec{r}_1) \quad , \tag{1}$$

where $\vec{r}$ is the separation of the centers of mass of the colliding target nucleus and the α particle, $\rho_T(\vec{r}_1)$ and $\rho_\alpha(\vec{r}_2)$ are the respective nucleon densities derived from nuclear charge distributions [12], and $t(E,\rho_T,\rho_\alpha,s)$ is the density-dependent effective NN interaction [1]. By means of the normalization factor λ the depth of the potential can be adjusted to elastic scattering data and to bound and resonant state energies of nuclear cluster states.

The imaginary part of the potential can be parametrized by a Woods-Saxon form.

The strength of the potential is measured by its volume integral per interacting pair of nucleons, e.g. for the real part

$$J_R(E) = \frac{4\pi}{A_P \cdot A_T} \int_0^\infty V(r,E) r^2 dr \quad , \tag{2}$$

Fig. 1a: Volume integrals of the real part of the optical potential for the α + ^{4}He, α + ^{16}O, and α + ^{40}Ca system.

Fig. 1b: Volume integrals of the imaginary part of the optical potential for the α + ^{4}He, α + ^{16}O, and α + ^{40}Ca system.

where A_P and A_T denote the projectile and target mass numbers, respectively.
In Figs. 1a and 1b the volume integrals for the real part, J_R, and the imaginary part, J_I, of the optical α-nucleus potential for the target nuclei ^{4}He, ^{16}O and ^{40}Ca are shown. These values have been obtained from the analysis of elastic scattering data and from calculations of ^{8}Be, ^{20}Ne and ^{44}Ti cluster states, respectively [3,4]. A strong energy and mass dependence for both, the real and imaginary parts, can be observed. The energy dependence of J_R (Fig. 1a) is due to the energy dependence of both, the effective NN interaction (Eq. 1) and the so-called dynamic polarization potential, which is related to that of the imaginary part of the potential by a dispersion relation [13,3]. The curves shown in Fig. 1a represent the results of calculations which consider both effects [4]. In these calculations the energy dependence of J_I was parametrized by [14]

$$\left.\begin{array}{ll} J_I = J_0 \dfrac{(E-E_0)^2}{(E-E_0)^2+\Delta^2} & E \geq E_0 \\ J_I = 0 & E < E_0 \end{array}\right\} \tag{3}$$

with E_0 being the threshold energy for inelastic processes. A linear regression procedure to the data points results in values for J_0 and Δ. The curves calculated with these parameters are given in Fig. 1b as solid, dashed and dotted lines.

3. Alpha-nucleus potentials for intermediate and heavy nuclei

Using the same procedure as described in Sec. 2, we determined optical alpha-nucleus potentials for intermediate and heavy nuclei. In Fig. 2 calculated elastic scattering cross sections for some target nuclei are compared with experimental data at different energies. An excellent agreement between the theoretical analysis and the experiment is obtained. From this fit to the experimental data, the strengths of the real part, λ, as well as the parameters of the imaginary Woods-Saxon potentials are obtained for energies above 15 MeV.

Fig. 2: Elastic α scattering on ^{90}Zr, ^{208}Pb, ^{70}Ge, ^{141}Pr and ^{144}Sm: Experimental data [25-37] and optical model fits calculated by using double-folded potentials.

Fig. 3a: Volume integrals of the real part of the optical α-nucleus potential for some intermediate and heavy nuclei.

Fig. 3b: Volume integrals of the imaginary part of the optical α-nucleus potential for some intermediate and heavy nuclei.

The volume integrals for the real as well as the imaginary part of these potentials are shown in Figs. 3a and 3b. The following results can be deduced:

(i) The mass dependence of the volume integrals for the real part of the potentials is very weak for all heavier nuclei. However, the absolute values are somewhat smaller than those for lighter nuclei (compare Figs. 1a and 3a).

(ii) The energy dependence has a similar form as the one obtained for lighter nuclei. The volume integrals of the real part have a maximum about $E_{CM} = 30$ MeV and decrease slightly when going to lower and higher energies. For astrophysically relevant energies, which are in the order of 10 MeV, a linear extrapolation to lower energies is performed (s. Fig. 3a). We obtain

$$J_R/(A_P \cdot A_T) = (320 + 0.67 E_{CM}) \text{ MeV fm}^3 \quad (E_{CM} \text{ in MeV}) \tag{4}$$

The curve shown in Fig. 3a is the result of a calculation for the α-^{208}Pb potential which contains the energy dependence of the folding potential and the dispersive part. The knowledge of bound and quasibound-state potentials is also necessary for the calculation of transfer and capture cross sections. We have calculated these alpha-nucleus potentials for some bound states. The volume integrals for these potentials are also close to the extrapolated values (s. Fig. 3a).

(iii) The volume integrals for the imaginary part of the potentials obtained from the fit to the experimental elastic scattering data for the target nuclei ^{70}Ge, ^{90}Zr, ^{141}Pr, ^{144}Sm and ^{208}Pb are shown in Fig. 3b. For ^{208}Pb the parametrization given in Eq. (3) was used to calculate the observed energy dependence of J_I (solid line in Fig. 3b). As expected a strong mass dependence of the volume integrals is observed, since the strength of the imaginary potential is quite different for a doubly magic and a strongly deformed nucleus. Therefore for an energy of $E_{CM} = 10$ MeV, J_I values can range between 10 and 60 MeV fm^3.

For unstable nuclei the mass densities necessary for the calculation of the folding potential cannot be obtained from electron scattering data. In these cases the densities can be calculated in the $(\sigma\omega\rho)$ model [15,16] used in the relativistic mean-field theory. We used the parameter set NLSH which is suited for neutron and proton rich nuclei [17,18,19]. We found that for stable tin and samarium isotopes the densities calculated in this model compare well with the experimental data [12]. In order to calculate the alpha-nucleus potentials for unstable nuclei, the strengths of the potentials were adjusted to reproduce the parametrized volume integrals given above.

4. Application to "p-process" isotopes

The alpha-nucleus potentials determined with the folding procedure are necessary for the calculation of (γ, α) photo disintegration cross sections in the p-process. As an example we consider the inverse reaction $^{144}\mathrm{Sm}(\alpha,\gamma)^{148}\mathrm{Gd}$. This reaction determines the ratio $^{142}\mathrm{Nd}/^{144}\mathrm{Nd}$ in some meteorites [20]. In previous work the astrophysical S-factors and photonuclear reaction rates have been generated using the statistical theory of nuclear reactions as employed by Michaud and Fowler [21]. They used an equivalent square well [ESW] deduced from a Woods-Saxon potential. However, the effective radius parameter for this ESW is quite uncertain [22]. For the two different ESW-radii for the α-^{144}Sm potential (R_α = 8.75 fm and 8.01 fm), the calculated cross sections differ by a factor of ten.

Table 1: S-factors and reaction rates for $^{144}\mathrm{Sm}(\alpha,\gamma)^{148}\mathrm{Gd}$

	ESW R = 8.75 fm	ESW R = 8.01 fm	Woods-Saxon potential [24]	Folding potential[a])
S-factor ($E_\alpha = 9.5$ MeV) [MeV · b]	$2.3 \cdot 10^{28}$ [b])	$2.3 \cdot 10^{27}$	$1.2 \cdot 10^{28}$	$7.8 \cdot 10^{27}$
reaction rate $T_9 = 2.5$ [cm^3 mol^{-1} s^{-1}]	$3.75 \cdot 10^{-15}$	$3.72 \cdot 10^{-16}$	$1.95 \cdot 10^{-15}$	$1.27 \cdot 10^{-15}$
reaction rate $T_9 = 3.0$ [cm^3 mol^{-1} s^{-1}]	$2.35 \cdot 10^{-12}$	$2.58 \cdot 10^{-13}$	$1.22 \cdot 10^{-12}$	$7.56 \cdot 10^{-13}$

[a]) this work

[b]) There is a misprint in [22] giving this value as $2.3 \cdot 10^{29}$.

In table 1 we list the astrophysical S-factor at the Gamow energy $E_\alpha = 9.5$ MeV and the reaction rates $T_9 = 2.5$ and $T_9 = 3.0$. In the first two columns the results of calculations using ESW [22] are given. In the third and fourth column the results of Hauser-Feshbach calculations using the code SMOKER [23] are shown. In the first case (column 3) an energy-independent Woods-Saxon potential with V = 185 MeV, W = 25 MeV, $R = R_W = 1.4 \cdot A^{1/3}$ fm and $a = a_W = 0.52$ fm [24], in the second case (column 4) for the real part a folding potential ($\lambda = 1.1573$) and for the imaginary part a Woods-Saxon potential ($W = 10$ MeV, $R_W = 1.4 \cdot A^{1/3}$ fm, $a_W = 0.52$ fm) was used.

With our improved folding potential the reaction rates shown in the last column of table 1 are about 1/3 of the value for an ESW with a radius of 8.75 fm. This corresponds about to 1/3 of the "recommended" value for the reaction rate in [22] giving a similar reaction rate as case C in table 1 of [22]. Therefore, our ^{146}Sm/^{144}Sm ratio is about 0.22 which is consistent with the cosmochemical data of 0.1 – 0.7 [20].

The astrophysical S-factor at $E_\alpha = 9.5$ MeV calculated in the direct-capture model gives an upper limit of about 10^{22} MeV $\cdot$ b, which is more than five orders of magnitude smaller than the Hauser-Feshbach result.

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